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Constraints on heavy colored scalars from Tevatron's Higgs exclusion limit

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Abstract

The null search for the Higgs boson at the Tevatron implies strong constraints on heavy colored particles that increase the gluon-fusion induced production rate of the Higgs. We investigate the implications of the Tevatron exclusion limit on example extensions of the Standard Model that contain a new scalar state transforming as either an adjoint or a fundamental under the QCD gauge group. The bounds on the adjoint (fundamental) scalar mass exceed 200 GeV (100 GeV) for natural choices of scalar-sector parameters.

1 Introduction

The hunt for the Higgs boson in order to uncover its role in electroweak symmetry breaking has been actively undertaken during the previous few years at the Tevatron accelerator complex. The CDF and D0 collaborations at the Tevatron have recently announced a 95% CL exclusion limit on a Standard Model (SM) Higgs boson with a mass in the ranges $100\,\text{GeV} \leq m_h \leq 109\,\text{GeV}$ and $158\,\text{GeV} \leq m_h \leq 175\,\text{GeV}$ [1]. This result is the culmination of intense efforts by experimentalists to reduce systematic errors inhibiting the search and to devise sophisticated multi-variate analysis techniques to enable control of the severe backgrounds [2], and by theorists to calculate precisely the production rate and distribution shapes of the Higgs within the SM [3]. Although the Tevatron will soon cease operation and the search for the Higgs boson will shift to the LHC, there is much to still be learned from Tevatron's null search. The lack of an observed Higgs signal, and our ability to calculate very precisely what is expected in the Standard Model, implies significant constrains on additional new states that increase the Higgs cross section. This has been demonstrated by CDF and D0 by their placement of stringent bounds on the existence of a fourth generation of fermions through its contribution to the gluon-fusion production of a Higgs [4]. In this work the collaborations placed model-independent bounds on $\sigma_{gg\to h} \times BR(h \to WW)$, allowing the implications of their null search on other SM extensions which affect the partonic process $gg \to h \to W^+W^-$ to be investigated.

Our goal in this manuscript is to further examine the implications of the Tevatron exclusion of the Higgs boson for physics beyond the Standard Model. We consider two example states that contribute significantly to the gluon-fusion production of a Higgs boson: color-fundamental and color-adjoint scalars. We calculate the gluon-fusion production cross section and W^+W^- decay width of the Higgs through next-to-next-to-leading-order (NNLO) in QCD using an effective-theory approach, and use the results of Ref. [4] to constraint the scalar parameter space. We have previously presented a detailed discussion of how color-adjoint scalars affect the Higgs production cross section [5], and we refer to this work for many technical details. For natural choices of scalar parameters, we find bounds of one hundred to several hundred GeV on the scalar mass, depending on both representation and the Higgs mass. While our interest in these specific particles is because of their utility as phenomenological examples, we note that such states appear in extensions of the Standard Models. For example, adjoint scalars in the $(\mathbf{8},\mathbf{1})_{\mathbf{0}}$ representation arise in theories with universal extra dimensions [6, 7]. We note that the examination of how the Tevatron exclusion limit impacts modes of new physics has received other attention in the literature [8].

Our manuscript is organized as follows. We review our calculational approach for determining the effects of the colored scalar on Higgs production in Section 2. A more detailed discussion is presented in our previous paper [5]. In Section 3 we present the Tevatron bounds on these states, derived using the model-independent constraints on Higgs production and decay via the partonic process $gg \to h \to W^+W^-$. We conclude in Section 4.

2 Calculational details

We discuss here our calculational procedure for obtaining the modifications in the Higgs production rate due to colored scalars. We consider two example models of colored states that can modify the SM prediction for gluon-fusion production of a Higgs boson: scalars in the $(\mathbf{8},\mathbf{1})_\mathbf{0}$ and $(\mathbf{3},\mathbf{1})_\mathbf{0}$ representations. The Lagrangians describing the interactions of these states with the SM are given by

$$\begin{aligned}\mathcal{L}^{adj} &= \mathcal{L}_{SM} + \mathrm{Tr}\left[D_\mu S D^\mu S\right] - m_S^{'2}\,\mathrm{Tr}\left[S^2\right] - g_s^2\,G_{4S}\,\mathrm{Tr}\left[S^2\right]^2 - \lambda_1 H^\dagger H\,\mathrm{Tr}\left[S^2\right],\\ \mathcal{L}^{fund} &= \mathcal{L}_{SM} + \left(D_\mu S\right)^\dagger D^\mu S - m_S^{'2}\,S^\dagger S - \frac{1}{2}g_s^2\,G_{4S}\,\left(S^\dagger S\right)^2 - \lambda_1 H^\dagger H\,S^\dagger S.\end{aligned} \tag{1}$$

In the adjoint Lagrangian $\mathcal{L}^{adj}$, S denotes the matrix-valued scalar field $S = S^A T^A$, while in the fundamental Lagrangian $\mathcal{L}^{fund}$, S denotes a vector in color space with three components. H indicates the Higgs doublet before electroweak symmetry breaking, v is the Higgs vacuum-expectation value, and D_μ is the covariant derivative. After electroweak symmetry breaking, the Higgs doublet is expanded as $H = \left(0, (v+h)/\sqrt{2}\right)$ in the unitary gauge. The masses of the colored scalars become $m_S^2 = m_S^{'2} + \lambda_1 v^2/2$. The Feynman rules which describe the scalar couplings to the Higgs boson h and to gluons are easily obtained from Eq. (1). The free parameters which govern the scalar properties are m_S, λ_1, and G_{4S}. We note that higher-order operators that break the $S \to -S$ symmetry present in the Lagrangians above, and which allow the scalar to decay, can be obtained in explicit models [6, 7]. We neglect them here since we anticipate that they have little effect on the $gg \to h$ production cross section. We also note that a quartic-scalar coupling is generated by QCD interactions even if it is set to zero at tree-level. At NNLO the quartic coupling must be included to obtain a renormalizable result, as demonstrated in our previous work [5]. We include this operator in the tree-level Lagrangian with a coefficient scaled by g_s^2, the QCD coupling constant squared, to permit an easier power-counting of loops.

We calculate the production rate of the Higgs through NNLO in QCD perturbation theory, which is necessary to reduce theoretical uncertainties arising from variations of the renormalization and facrorization scales in this process. We perform this calculation using an effective-theory strictly valid when the Higgs mass is less than twice the scalar mass, and also less than twice the top mass, $m_h < 2m_{S,t}$. In this limit, both the top quark and colored scalar can be integrated out to obtain the following effective Lagrangian:

$$\mathcal{L}^{eff} = \mathcal{L}_{QCD}^{n_l,eff} - C_1\,\frac{H}{v}\,\mathcal{O}_1, \tag{2}$$

where C_1 is a Wilson coefficient and the operator $\mathcal{O}_1$ is

$$\mathcal{O}_1 = \frac{1}{4}\,G_{\mu\nu}^a G^{a\mu\nu}\,. \tag{3}$$

This form is valid for both the adjoint and fundamental scalar; only C_1 changes. This effective-theory has been extensively used to determine the contribution of top-quark loops to the SM Higgs-gluon coupling, and to calculate the gluon-fusion cross section for Higgs production at hadron colliders [9–18]. It has also been used to determine the effect of a

fourth generation of fermions on the production rate at the Tevatron [19]. When normalized to the full m_t-dependent leading-order result, the effective-theory reproduces the exact NLO result obtained in Ref. [12,20] to better than 1% for $m_h < 2m_t$ and to 10% or better for Higgs boson masses up to 1 TeV. For scalars in the adjoint representation, this approximation has been studied against the exact NLO calculation [21], and is again accurate to the $1-2\%$ percent level for $m_h \leq 2m_S$. The deviation reaches a maximum of 10% for Higgs masses much heavier than the scalar, except very near the threshold $m_h \approx 2m_S$. We adopt this normalized effective-theory as the basis for our analysis. A detailed derivation of the Wilson coefficient for the adjoint scalar was presented in Ref. [5]. We have derived C_1 for the fundamental scalar, and present it in Eq. (6) of the Appendix.

Before presenting numerical results, we explain what contributions we include in the gluon-fusion cross section and gluinic decay width of the Higgs, both of which are modified in the presence of the scalar. The leading-order amplitude for the $gg \to h$ process takes the form

$$\mathcal{A}^{LO} = \mathcal{A}_t^{LO} + \mathcal{A}_b^{LO} + \mathcal{A}_S^{LO}, \tag{4}$$

where the subscripts t, b, S respectively denote the top, bottom, and scalar contributions. The amplitude for $h \to gg$ has the identical structure. Upon squaring this amplitude, interferences between the contributions of each particle are obtained. We denote by σ_{t+S} the terms obtained by squaring together the top and scalar amplitudes, and keeping both the interference term and the pieces from each separate particle squared. We let σ_{tb}, σ_{Sb} denote the interferences between the bottom-quark amplitude with the top and the scalar pieces. For the cross section at the n-th order in perturbation theory, we use the following expression:

$$\sigma^n = \sigma_{t+S}^{LO}(m_t, m_S)\, K_{EFT}^n + \sigma_{Sb}^{LO}(m_S, m_b) + \sigma_{tb}^{LO}(m_t, m_b) + \sigma_{bb}^{LO}(m_b). \tag{5}$$

K_{EFT}^n denote the ratio of the n-th order result for the cross section over the LO result, with both quantities computed in the effective-theory defined in Eq. (2). The expression multiplying the K-factors is the LO cross section maintaining the exact dependence on the scalar and top-quark masses. The remaining terms account for the scalar-bottom interference, the top-bottom interference, and the bottom-squared contribution at LO with their exact mass dependences. Various electroweak corrections which modify the SM contribution at the percent level [22–26] are not known for the scalar, and for consistency are neglected. As discussed above, this calculational framework yields results usually accurate to the percent level, and to 10% at worst. For the SM it gives results differing by only a couple of percent from the official predictions utilized by the Tevatron collaborations [24, 27], and also closely matches the prescription for LHC cross sections adopted by ATLAS and CMS [28]. To calculate the partial width of the Higgs boson into gluons, we use HDECAY [29] to calculate its SM partial width. We then scale its result by the ratio of the amplitude in Eq. (4) squared over that in the SM squared, together with a factor accounting for the different Wilson coefficients in the SM and in the presence of the scalar. This partial width is then used with the other outputs of HDECAY to form the $h \to W^+W^-$ branching ratio.

3 Numerical results

We now explore what regions of scalar parameter space are excluded by the model-independent search of the CDF and D0 collaborations for the process $gg \to h \to W^+W^-$. We study the adjoint and fundamental representations separately. The addition of a scalar to the spectrum induces two competing effects on the Tevatron signal. The scalar tends to increase the gluon-fusion production cross section throughout the parameter space. However, it also increases the partial decay width of the Higgs into gluons, which decreases the branching ratio into W^+W^-. To derive the allowed region of scalar parameter space we calculate $\sigma_{gg\to h} \times BR(h \to WW)$, and demand that it be less than the Tevatron limit on this quantity from Ref. [4]. In addition, to avoid strong couplings and a breakdown of our perturbative analysis, we impose the additional constraint $\Gamma_{total}/m_h < 1/5$ on the total width of the Higgs. This constraint prevents the ratio $\lambda_1 v^2/m_S^2$ appearing in the Wilson coefficient, where v is the vacuum-expectation value of the Higgs, from becoming too large.

We use MSTW parton distribution functions [30] extracted through NNLO in QCD when evaluating the cross section in Eq. 5. For the bottom quark we use the pole mass $m_b = 4.75\,\text{GeV}$; this leads to slightly more conservative bounds than those obtained with the $\overline{MS}$ mass. The cross section is evaluated using the scale choice $\mu_R = \mu_F = m_h/2$. In the scalar sector, we must set the parameters λ_1, G_{4S}, and m_S. To determine the reasonable range of G_{4S}, we perform a renormalization-group analysis and demand that this coupling does not encounter a Landau pole until a cutoff $\Lambda = 10\,\text{TeV}$. This was discussed in detail in Ref. [5], to which we refer the reader for further details. For the adjoint scalar, this leads to the approximate condition $G_{4S}(v) < 1.5$, while for the fundamental we find $G_{4S}(v) < 2.5$. The value of G_{4S} does not significantly modify the bounds we derive. While we set $G_{4S}(v) = 1$ in our study here, we have checked that other values in the allowed region do not modify our bounds by more than 5%. The cross section depends primarily on λ_1 through the ratio λ_1/m_S^2, as discussed in Ref. [5] . The constraints on m_S depend strongly on the value of λ_1 chosen. However, there is no symmetry reason to expect a small value for this coupling, indicating a 'natural' value $\lambda_1 \approx 1$. We note that since the scalars considered here will not significantly alter predictions for the precision electroweak observables measured at LEP and SLC, the Higgs mass is also expected to be in the approximate range $m_h < 200$ GeV. Results for other values λ_1^o, but keeping the mass fixed at m_S, can be approximately obtained by studying the results at a mass m_S^o given by $\lambda_1/m_S^2 = \lambda_1^o/(m_S^o)^2$. To avoid three-dimensional plots we simply set $\lambda_1 = 1$ in our analysis, and note that results for other values can be obtained by scaling $\sigma_{gg\to h} \times BR(h \to WW)$ as indicated. The only remaining parameters are m_h and m_S. We present our results as exclusion regions in this two-dimensional space.

The excluded regions for the adjoint and fundamental scalars are presented in Figs. 1 and 2, respectively. We note that these are 95% CL exclusion bounds. The strongest bounds occur when $m_h = 165\,\text{GeV}$ in both cases. Since this search utilizes only the gluon-fusion production mode, the SM Higgs is not excluded in this analysis, unlike in the combined global analysis [1]. Taking $m_S \to \infty$ gives the SM result for the production rate, and if the SM was excluded for a certain Higgs mass range in this analysis, even an extremely heavy scalar would not be allowed. Adjoint-scalar masses approaching 1 TeV are excluded for this Higgs mass, while fundamental-scalar masses near 500 GeV are ruled out. The constraints are considerably weaker for other values of m_h, but adjoint-scalar masses less than 130 GeV

are ruled out for Higgs masses between 135 GeV and 250 GeV, while fundamental-scalar masses less than 100 GeV are excluded for Higgs masses from 150 GeV to 190 GeV. For comparison, the direct-search limits on the adjoint scalar mass is estimated to be 280 GeV when the scalar decays primarily into $b\bar{b}$ at the Tevatron [7]. The direct search constraint is very sensitive to the decay mode of the scalar. However, it is also insensitive to m_h and λ_1; the two techniques for probing the scalar parameter space are complementary. The regions of light scalar mass are allowed because $Br(h \to WW)$ drops quickly as m_S is decreased. However, scalar masses that are too light increase the gluonic partial width to the point that Γ_{total}/m_h becomes too large, leading to the exclusion at small values of m_S. The exclusion regions extending to large m_S apparent in the figures regions at high m_h arise from the increase of the cross section near the threshold $m_h \approx 2m_S$. The allowed values of scalar masses found by our analysis are presented in Tables 1 and 2. We note that we consider only the Tevatron Higgs exclusion limit in determing these ranges. Also, the very light allowed masses near $m_S \approx 10$ GeV should not be taken too seriously due to the limitations of the effective-theory analysis. While caution must be exercised in using our effective-theory framework for very light scalar masses, we note that the deviation of Eq. (5) form the exact next-to-leading order calculation of Ref. [21] was found to be less than 10% for $m_h/m_S \leq 5$.

We note that the bulk of the exclusion region for the adjoint scalar with a mass in the range $120 - 250$ GeV and the fundamental scalar with a mass in the range $130 - 200$ GeV comes from the Tevatron bounds. The finite width constraint has no effect in these regions. Also, in this dominant part of the excluded parameter space the effective-theory is working with high precision. The properties of the width and the effective-theory limitations come in only in the tails extending to $m_h \geq 200$ GeV and very light scalar masses, where the Tevatron cross section becomes too small to constrain the parameter space.

As a final comment, we also note that the validity of the Tevatron exclusion limits have been actively debated within the past year [31], due to the collaborations' treatment of theoretical systematic errors. We adopt the official results of the CDF and D0 collaborations in our analysis.

4 Conclusions

In this manuscript we have investigated the implications that Tevatron's null search for the process $gg \to h \to W^+W^-$ has for extensions of the Standard Model. We have considered two example particles that substantially alter the gluon-fusion production of the Higgs and its branching fraction into W^+W^-, and that also exist in extensions of the SM: heavy scalars in the adjoint and fundamental representations of the color gauge group. We have outlined an effective-theory computation of how the colored scalars modify Higgs properties, and have also presented the Wilson coefficient for the fundamental scalar needed for studies of this state through NNLO in QCD perturbation theory. Using the model-independent bounds on $\sigma_{gg\to h} \times BR(h \to WW)$ presented in Ref. [4], we have derived the excluded regions of scalar parameter space for both example states. The constraints can be quite severe; for Higgs masses near $m_h \approx 165$ GeV, adjoint scalar masses approaching 1 TeV are ruled out, while fundamental scalar masses near 500 GeV are not allowed. These should be compared to the estimated direct search reach of $m_S \approx 280$ GeV for the adjoint scalar [7]. Of course, the

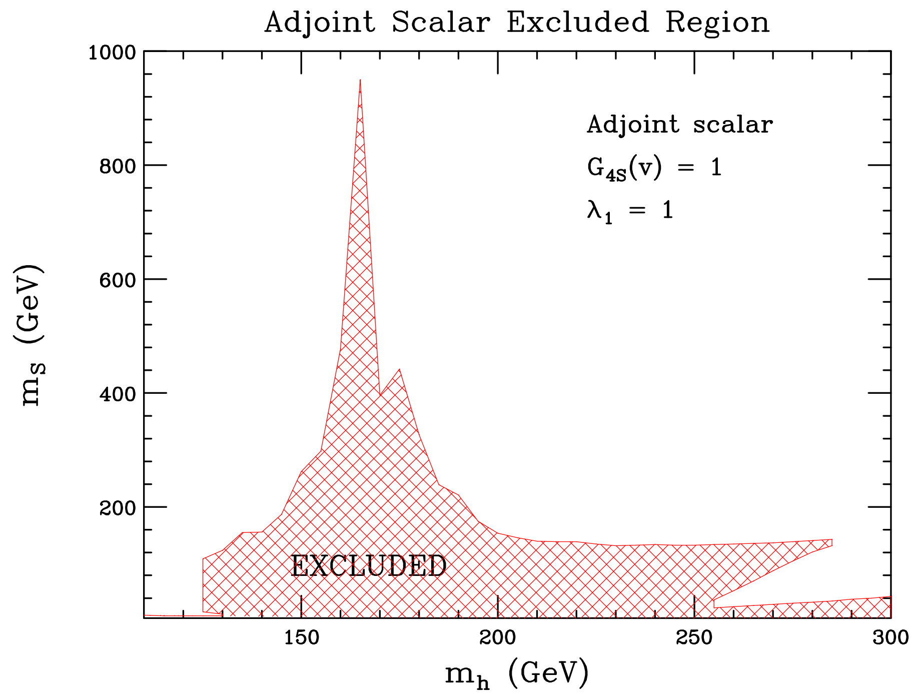

Figure 1: Bounds on the (m_h, m_S) parameter space for an adjoint scalar. The excluded values are denoted by red hatching. The choices of $G_{4S}(v)$ and λ_1 are discussed in the text.

indirect bounds here depend significantly on m_h and λ_1, which the direct-search limit does not. However, it is also independent of the scalar decay, which plays a crucial role in the direct search. Throughout the natural range of m_h and λ_1, the Tevatron results restrict the colored-scalar mass to be greater than $100-200$ GeV, with the stronger restriction occurring for the adjoint representation.

The Tevatron null search for the Higgs boson has implications beyond limiting the Standard Model parameter space. It imposes severe constraints on new theories that contain heavy colored states which modify the Higgs-gluon interaction. We have demonstrated this by showing the strong limits obtainable on color adjoint and fundamental scalars. It would be interesting to investigate other consequences of the Tevatron result for physics beyond the Standard Model.

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Appendix

We present here the Wilson coefficient C_1 for the fundamental scalar that appears in the

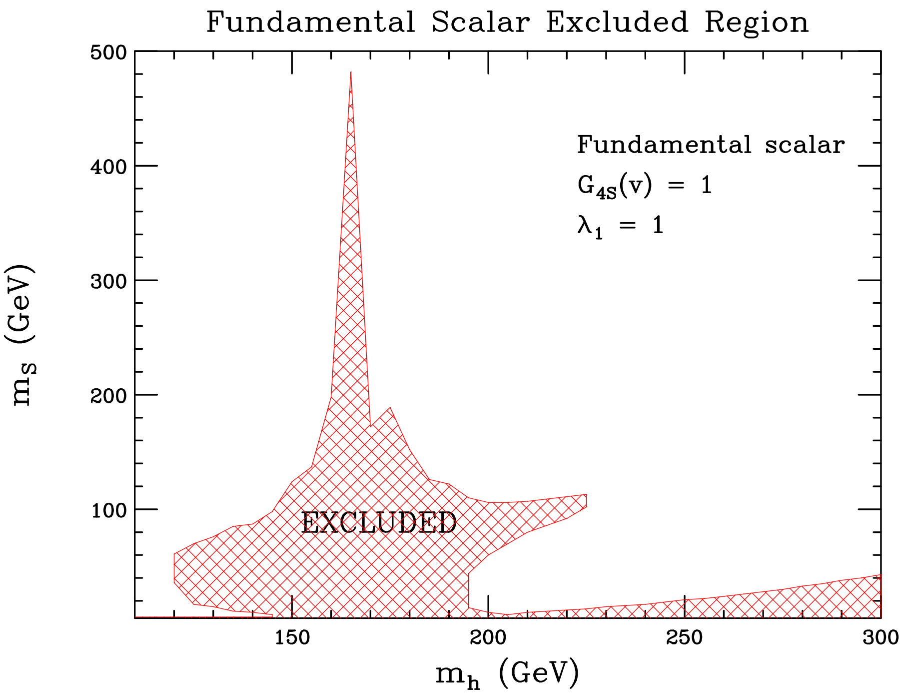

Figure 2: Bounds on the (m_h, m_S) parameter space for a fundamental scalar. The excluded values are denoted by red hatching. The choices of $G_{4S}(v)$ and λ_1 are discussed in the text.

effective Lagrangian of Eq. (2):

$$
\begin{aligned}
C_1 = \; & a\left[-\frac{\lambda_1 v^2}{48\, m_S^2} - \frac{1}{3}\right] \;+\; a^2\left[\frac{\lambda_1 v^2}{16\, m_S^2}\left(-\frac{3}{2} + \frac{G_{4S}}{3}\right) - \frac{11}{12}\right] \\
& + \; a^3\left[-\frac{(225x^6 - 72x^4 + 131x^2 - 228)\ln^2(x)}{6144(x-1)x^2(x+1)} + \frac{(861x^4 + 7435x^2 + 684)\ln(x)}{9216\, x^2}\right. \\
& + \frac{4608\, L_S\, x^4 - 105x^4 + 21888\, L_S\, x^2 - 89129x^2 - 2052}{27648\, x^2} + n_l\left(\frac{1}{288}(64\, L_S + 67) + \frac{2\ln(x)}{9}\right) \\
& - \frac{\lambda_1 v^2}{2\, m_S^2}\left\{\frac{7}{96}(2\, L_S - 1)\, G_{4S}^2 + \left(\frac{424\, L_S - 823}{2304} + \frac{\ln(x)}{72}\right) G_{4S} + \frac{(-6\, L_S - 71)\, n_l}{1728}\right. \\
& - \frac{(225x^6 - 72x^4 + 131x^2 - 228)\ln^2(x)}{12288(x-1)x^2(x+1)} + \frac{(287x^4 - 103x^2 + 228)\ln(x)}{6144\, x^2} \\
& \left. + \frac{4608\, L_S\, x^4 - 2409x^4 - 25152\, L_S\, x^2 + 48247x^2 - 2052}{55296\, x^2}\right\} + \left(75x^6 + 26x^4 + 7x^2 + 76\right) \times \\
& \left((\ln(1+x) - \ln(1-x))\frac{\ln^2(x)}{4096\, x^3} + (\mathrm{Li}_2(-x) - \mathrm{Li}_2(x))\frac{\ln(x)}{2048\, x^3} + \frac{\mathrm{Li}_3(x) - \mathrm{Li}_3(-x)}{2048\, x^3}\right. \\
& \left.\left. - \frac{\lambda_1 v^2}{2\, m_S^2}\left\{(\ln(1+x) - \ln(1-x))\frac{\ln^2(x)}{8192\, x^3} + (\mathrm{Li}_2(-x) - \mathrm{Li}_2(x))\frac{\ln(x)}{4096\, x^3} + \frac{\mathrm{Li}_3(x) - \mathrm{Li}_3(-x)}{4096\, x^3}\right\}\right)\right].
\end{aligned}
\tag{6}
$$

In these expressions, we have set $L_i = \ln(m_i/\mu)$ for $i = (S, t)$, $x = m_t/m_S$ and $a = g_s^2/(4\pi^2)$. As discussed in our previous work [5], the combination $\lambda_1 v^2/m_S^2$ is scale-invariant if only

m_h (GeV)	m_S (GeV)	m_h (GeV)	m_S (GeV)
110	$10-\infty$	210	$140-\infty$
115	$9-\infty$	215	$139-\infty$
120	$9-\infty$	220	$139-\infty$
125	$9-16$, $109-\infty$	225	$135-\infty$
130	$9-12$, $124-\infty$	230	$132-\infty$
135	$155-\infty$	235	$133-\infty$
140	$156-\infty$	240	$134-\infty$
145	$187-\infty$	245	$133-\infty$
150	$262-\infty$	250	$133-\infty$
155	$298-\infty$	255	$23-37$, $134-\infty$
160	$477-\infty$	260	$25-53$, $135-\infty$
165	$950-\infty$	265	$27-70$, $136-\infty$
170	$396-\infty$	270	$29-88$, $137-\infty$
175	$442-\infty$	275	$31-105$, $139-\infty$
180	$327-\infty$	280	$33-121$, $141-\infty$
185	$239-\infty$	285	$35-132$, $143-\infty$
190	$221-\infty$	290	$38-\infty$
195	$175-\infty$	295	$40-\infty$
200	$154-\infty$	300	$43-\infty$
205	$146-\infty$	–	–

Table 1: Allowed values of the adjoint scalar mass m_S for each m_h, given the Tevatron bounds [4] and the constraint $\Gamma_{total}/m_h < 1/5$. The choices of the parameters λ_1 and G_{4S} are described in the text.

QCD-induced α_S corrections are considered. The Wilson coefficient therefore takes the same form in both the $\overline{MS}$ and pole schemes through next-to-next-to-leading order in perturbation theory. In presenting our numerical results we interpret the masses of the top and scalar as pole masses.

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m_h (GeV)	m_S (GeV)	m_h (GeV)	m_S (GeV)
110	$6-\infty$	210	$10-80$, $107-\infty$
115	$6-\infty$	215	$11-86$, $109-\infty$
120	$6-36$, $61-\infty$	220	$12-92$, $111-\infty$
125	$6-17$, $70-\infty$	225	$13-102$, $113-\infty$
130	$6-15$, $76-\infty$	230	$15-\infty$
135	$6-11$, $85-\infty$	235	$16-\infty$
140	$6-10$, $87-\infty$	240	$17-\infty$
145	$6-8$, $98-\infty$	245	$19-\infty$
150	$124-\infty$	250	$21-\infty$
155	$137-\infty$	255	$22-\infty$
160	$198-\infty$	260	$24-\infty$
165	$482-\infty$	265	$26-\infty$
170	$172-\infty$	270	$28-\infty$
175	$189-\infty$	275	$30-\infty$
180	$152-\infty$	280	$33-\infty$
185	$126-\infty$	285	$35-\infty$
190	$122-\infty$	290	$38-\infty$
195	$14-44$, $110-\infty$	295	$40-\infty$
200	$10-60$, $106-\infty$	300	$43-\infty$
205	$9-70$, $106-\infty$	–	–

Table 2: Allowed values of the fundamental scalar mass m_S for each m_h, given the Tevatron bounds [4] and the constraint $\Gamma_{total}/m_h < 1/5$. The choices of the parameters λ_1 and G_{4S} are described in the text.

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